

Z' Mediated Flavor Changing Neutral Currents in B Meson DecaysVernon Barger,^{1,*} Cheng-Wei Chiang,^{1,†} Paul Langacker,^{2,‡} and Hye-Sung Lee^{1,§}¹*Department of Physics, University of Wisconsin, Madison, WI 53706*²*Department of Physics and Astronomy, University of Pennsylvania, Philadelphia, PA 19104*

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We study the effects of an extra $U(1)'$ gauge boson with flavor changing couplings with fermion mass eigenstates on certain B meson decays that are sensitive to such new physics contributions. In particular, we examine to what extent the current data on $B_d \rightarrow \phi K$, $B_d \rightarrow \eta' K$ and $B_s \rightarrow \mu^+ \mu^-$ decays may be explained in such models, concentrating on the example in which the flavor changing couplings are left-chiral. We find that within reasonable ranges of parameters, the Z' contribution can readily account for the anomaly in $S_{\phi K_S}$ but is not sufficient to explain large branching ratio of $B_d \rightarrow \eta' K$ with the same parameter value. $S_{\phi K_S}$ and $S_{\eta' K_S}$ are seen to be the dominant observables that constrain the extra weak phase in the model.

I. INTRODUCTION

CP violation has been a puzzling phenomenon in the studies of elementary particle physics since the first observation of its effects in hadronic kaon decays almost four decades ago [1]. In the standard model (SM), CP -violation is due entirely to the Cabibbo-Kobayashi-Maskawa (CKM) mechanism [2, 3], describing the mismatch between the unitary transformations relating the up and down type quark mass eigenstates to the corresponding weak eigenstates. The CKM matrix involves a single weak phase along with three mixing angles. The validity of the CKM picture is further strengthened by the fact that recent $\sin 2\beta$ measurements from time-dependent CP asymmetries of decay modes involving the $b \rightarrow c\bar{c}s$ subprocess [4] agree well with the range of the weak phase β from many other constraints [5]. However, it is still unknown whether there are any other sources that may give rise to CP -violating effects. Good places to search for deviations from the SM predictions are decay processes that are expected to be rare in the SM, which may reveal new physics through interference effects. In particular, discrepancies among the time-dependent CP asymmetries of different B decay modes may show evidence for new physics [6, 7, 8, 9, 10, 11].

Recently, an anomaly was reported in the time-dependent CP asymmetry measurement of the $B_d \rightarrow \phi K_S$ decay mode. Within the framework of the SM, this process should also provide us with information on the weak phase β , up to about 5% theoretical uncertainty [6, 12]. However, the averaged value of $S_{\phi K_S}$ reported by the BaBar and Belle groups is [4] (individual measurements to be quoted in Table II)

$$S_{\phi K_S} = -0.147 \pm 0.697 \quad (S = 2.11) . \quad (1)$$

This result is only about 1.3σ away from the corresponding quantity measured by the $B \rightarrow J/\psi K_S$ mode, $S_{J/\psi K_S} = 0.736 \pm 0.049$ [4]. However, the scale factor $S = 2.11$ suggests a discrepancy between the two experimental groups [57]. Before this discrepancy is settled, the difference between $S_{\phi K_S}$ and $S_{J/\psi K_S}$ suggests the possibility of new physics contributions. From the theoretical point of view, the $B \rightarrow \phi K_S$ decay is a loop-induced process involving $b \rightarrow s\bar{s}s$ penguin operators in the SM. Therefore, it is susceptible to new physics contributions even if they are suppressed by a large mass parameter which characterizes the new physics scale.

In addition to model-independent approaches [6, 12, 14], many studies have been made to explain the anomaly in supersymmetric and related models [15, 16]. Such an effect can also be explained using models in which the bottom quark is mixed with heavy mirror fermions with masses of the order of the weak scale [17]. It is the purpose of this work to show that a new physics effect of similar size can be obtained from some models with an extra Z' boson.

Z' bosons are known to naturally exist in well-motivated extensions of the SM [18]. The Z' mass is constrained by direct searches at Fermilab, weak neutral current data, and precision studies at LEP and the SLC [19, 20, 21], which

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give a model-dependent lower bound around 500 GeV. The latter also severely limits the Z - Z' mixing angle $|\theta| < \text{a few} \times 10^{-3}$. A Z' could be relevant to the NuTeV experiment [22] and, if the couplings are not family universal [20, 23], to the anomalous value of the forward-backward asymmetry A_{FB}^b [21]. (Earlier hints of a discrepancy in atomic parity violation have largely disappeared due to improved calculations of radiative corrections [24].) We therefore study the Z' boson in the mass range of a few hundred GeV to 1 TeV, assuming no mixing between Z and Z' .

Interesting phenomena arise when the Z' couplings to physical fermion eigenstates are non-diagonal. This is possible if there exist additional exotic fermions that have different $U(1)'$ charges from the ordinary fermions, as found in E_6 models [25, 26, 27]. However, in these models left-handed fermion mixings may induce undesirable flavor changing neutral currents (FCNC) mediated by the Z boson even in the absence of Z - Z' mixing or nonuniversal family couplings. One can avoid this consequence by confining the mixing to be between right-handed fermions and the exotic quarks [28]. Alternatively, other models give family nonuniversal Z' couplings as a result of different ways of constructing families in some string models [29, 30, 31, 32]. FCNC and possibly new CP -violating phenomena will also occur in these models after fermion mixings are taken into account. These can occur for both left and right-handed fermions.

Although experiments on FCNC processes (such as the mass difference between K_L and K_S and the $K_L \rightarrow \mu^+ \mu^-$ decay) have significantly constrained the Z' couplings of the first and second generation quarks to be almost the same and diagonal, the couplings to the third generation quarks are not well constrained. Similar statements apply to the charged leptons. It has been shown in Ref. [29, 30, 31, 32] that indeed the third generation fermions can have different Z' couplings from the other two generations.

We use all of the above-mentioned features to study the imprints of the Z' boson on certain processes that involve $b \rightarrow s$ transitions. In Section II, we present the model and framework to be studied. In Section III, we show the constraints on the model parameters from the current data of $S_{\phi K_S}$, $A_{\phi K_S}$ and the branching ratio $\mathcal{B}(B_d \rightarrow \phi K)$. In Section IV, we study a related process $B_d \rightarrow \eta' K_S$, also including its CP asymmetries and branching ratio. Another way to constrain the b - s - Z' coupling is the leptonic decay process $B_s \rightarrow \mu^+ \mu^-$. We study this process in Section V and show how sensitively Run II at Fermilab can probe this coupling in the near future. We conclude in Section VI.

II. FORMALISM

In this paper, we concentrate on models in which the interactions between the Z' boson and fermions are flavor nonuniversal for left-handed couplings and flavor diagonal for right-handed couplings. The analysis can be straightforwardly extended to general cases in which the right-handed couplings are also nonuniversal across generations. The basic formalism of flavor changing effects in the Z' model with family nonuniversal and/or nondiagonal couplings has been laid out in Ref. [23], to which we refer readers for detail. Here we just briefly review the ingredients needed in this paper.

We write the Z' term of the neutral-current Lagrangian in the gauge basis as

$$\mathcal{L}^{Z'} = -g' J'_\mu Z'^\mu, \quad (2)$$

where g' is the gauge coupling associated with the $U(1)'$ group at the M_W scale. We neglect its renormalization group (RG) running between M_W and $M_{Z'}$. The Z' boson is assumed to have no mixing with the SM Z boson [58]. The chiral current is

$$J'_\mu = \sum_{i,j} \bar{\psi}_i^I \gamma_\mu [(\epsilon_{\psi_L})_{ij} P_L + (\epsilon_{\psi_R})_{ij} P_R] \psi_j^I, \quad (3)$$

where the sum extends over the flavors of fermion fields, the chirality projection operators are $P_{L,R} \equiv (1 \mp \gamma_5)/2$, the superscript I refers to the gauge interaction eigenstates, and ϵ_{ψ_L} (ϵ_{ψ_R}) denote the left-handed (right-handed) chiral couplings. ϵ_{ψ_L} and ϵ_{ψ_R} are hermitian under the requirement of a real Lagrangian. The fermion Yukawa coupling matrices Y_ψ in the weak basis can be diagonalized as

$$Y_\psi^D = V_{\psi_R} Y_\psi V_{\psi_L}^\dagger \quad (4)$$

using the bi-unitary matrices $V_{\psi_{L,R}}$ in $\psi_{L,R} = V_{\psi_{L,R}} \psi_{L,R}^I$, where $\psi_{L,R}^I \equiv P_{L,R} \psi^I$ and $\psi_{L,R}$ are the mass eigenstate fields. The usual CKM matrix is then given by

$$V_{\text{CKM}} = V_{u_L} V_{d_L}^\dagger. \quad (5)$$

The chiral Z' coupling matrices in the physical basis of down-type quarks thus read

$$B_d^L \equiv V_{d_L} \epsilon_{d_L} V_{d_L}^\dagger, \quad (6)$$

$$B_d^R \equiv V_{d_R} \epsilon_{d_R} V_{d_R}^\dagger, \quad (7)$$

where the $B_d^{L,R}$ are hermitian. We do not need the corresponding couplings for up-type quarks or charged leptons in our discussions.

As long as $\epsilon_{dL,R}$ is not proportional to the identity matrix, $B_d^{L,R}$ will have nonzero off-diagonal elements that induce FCNC interactions. To see this, consider as an example the simplified ϵ_{dL} matrix for the down-type quarks of the form

$$\epsilon_{dL} = Q_d \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & X \end{pmatrix}, \quad (8)$$

where both d and s quarks have the same Z' charge Q_d and X is the ratio of the Z' charge of b to Q_d . If we assume the mixing is among the down-type quarks only, $V_{dL}^\dagger = V_{\text{CKM}}$ and $V_{uL} = \mathbf{1}$. Any redefinition of the quark fields by pure phase shifts would have no effect on the resultant B_d^L . All the off-diagonal couplings are proportional to the Z' charge difference between the b quark and the d, s quarks, as expected. Using the standard parametrization [33], the explicit form of the off-diagonal Z' coupling between b and s quarks, for example, is

$$B_{sb}^L = (1 - X)Q_d \cos \theta_{13} \cos \theta_{23} (\cos \theta_{12} \sin \theta_{23} + \cos \theta_{23} \sin \theta_{12} \sin \theta_{13} e^{-i\delta_{13}}), \quad (9)$$

where θ_{ij} are the mixing angles between the i th and the j th generations and δ_{13} is the CP -violating weak phase. In this example, B_{sb}^L is proportional to the product $V_{tb}V_{ts}^*$ of elements of the CKM matrix.

More generally, one can always pick a basis for the weak eigenstates in which the $\epsilon_{dL,R}$ matrices are diagonal and of the form (8), though with different Q_d and X for the ϵ_{dL} and ϵ_{dR} . However, the Yukawa matrices Y_ψ^u and Y_ψ^d will in general not be diagonal in that basis, so that $V_{uL} \neq \mathbf{1}$ and $V_{dL}^\dagger \neq V_{\text{CKM}}$. In that case $B_d^{L,R} = V_{dL,R} \epsilon_{dL,R} V_{dL,R}^\dagger$ will in general be nondiagonal and complex, with new mixing angles and CP violating phases not directly related to V_{CKM} [59].

Instead of restricting ourselves to models with particular parameter choices in the couplings and mixings, we will take the effective theory point of view and constrain the effective couplings relevant to the decay modes of interest in the following analysis. However, to be more definite, we assume that the right-handed coupling matrix B_d^R is flavor diagonal. If B_d^R is nondiagonal, new operators involving different chirality structures will be induced in B decays.

III. $B_d \rightarrow \phi K_S$

Within the SM, the $B^0 \rightarrow \phi K^0$ decay proceeds through the loop-induced $b \rightarrow s\bar{s}s$ transition, which involves dominantly the QCD penguin but also some electroweak (EW) and chromomagnetic penguin contributions. To illustrate possible modifications due to the existence of an extra $U(1)'$ gauge boson, we will neglect the smaller contributions from weak annihilation diagram in the following analysis although they can play some role in enhancing the branching ratios [34]. This two-body hadronic B meson decay can be conveniently analyzed in the framework of the effective weak Hamiltonian and factorization formalism [35, 36].

Since the penguin diagrams receive dominant contributions from the top quark running in the loop, the effective Hamiltonian relevant for the charmless $|\Delta S| = 1$ decays can be written as

$$\mathcal{H}_{\text{eff}}^{\text{SM}} = \frac{G_F}{\sqrt{2}} \left\{ V_{ub}V_{us}^* [c_1 O_1 + c_2 O_2] - V_{tb}V_{ts}^* \left[\sum_{i=3}^{10} c_i O_i + c_g O_g \right] \right\} + \text{h.c.} \quad (10)$$

Here

$$O_1 = (\bar{u}_\alpha b_\alpha)_{V-A} (\bar{s}_\beta u_\beta)_{V-A}, \quad O_2 = (\bar{u}_\alpha b_\beta)_{V-A} (\bar{s}_\beta u_\alpha)_{V-A} \quad (11)$$

are tree-level color-favored and color-suppressed operators,

$$O_{3(5)} = (\bar{s}_\alpha b_\alpha)_{V-A} (\bar{s}_\beta s_\beta)_{V-A(V+A)}, \quad O_{4(6)} = (\bar{s}_\alpha b_\beta)_{V-A} (\bar{s}_\beta s_\alpha)_{V-A(V+A)} \quad (12)$$

are the QCD penguin operators,

$$O_{7(9)} = \frac{3}{2} e_s (\bar{s}_\alpha b_\alpha)_{V-A} (\bar{s}_\beta s_\beta)_{V+A(V-A)}, \quad O_{8(10)} = \frac{3}{2} e_s (\bar{s}_\alpha b_\beta)_{V-A} (\bar{s}_\beta s_\alpha)_{V+A(V-A)} \quad (13)$$

are the EW penguin operators ($e_s = -1/3$ is the electric charge of the strange quark), and

$$O_g = \frac{g_s}{8\pi^2} m_b \bar{s}_\alpha \sigma^{\mu\nu} T_{\alpha\beta}^a (1 + \gamma_5) b_\beta G_{\mu\nu}^a \quad (14)$$

TABLE I: The SM Wilson coefficients used in the present analysis. We assume the naive factorization for a_i (i.e. $N_C^{\text{eff}} = 3$), and ignore small differences between the $b \rightarrow s$ and $\bar{b} \rightarrow \bar{s}$ decays, expecting more significant effects from new physics. c_i^{eff} and a_i ($i = 3 \cdots 10$) should be multiplied by 10^{-5} .

Operator	c_i^{eff}	$a_i(N_C^{\text{eff}} = 3)$
O_1	1.198	1.064
O_2	-0.403	-0.004
O_3	$2817 + 301i$	815
O_4	$-6006 - 903i$	$-5067 - 803i$
O_5	$2036 + 301i$	-425
O_6	$-7384 - 903i$	$-6705 - 803i$
O_7	$-28 - 12i$	$-5 - 12i$
O_8	70	$60 - 4i$
O_9	$-1079 - 12i$	$-957 - 12i$
O_{10}	366	$6 - 4i$

is the chromomagnetic operator, where $(\bar{q}_1 q_2)_{V \pm A} \equiv \bar{q}_1 \gamma_\mu (1 \pm \gamma_5) q_2$ and α, β refer to color indices.

We mention in passing that the Z' boson will also modify the $|\Delta B| = 2$ effective Hamiltonian relevant to $B_d - \bar{B}_d$ mixing, but in an unnoticeable way. This is because the additional contribution is proportional to the square of the Z' couplings between the first and third generations, $|B_{db}^{L,R}|^2$, which is much more suppressed than the SM contribution. Although the Z' also contributes to $b \rightarrow (c\bar{c})s$ transitions at the tree level and gains a color factor relative to the SM tree process, it is nevertheless suppressed by the $B_{sb}^{L,R}$ couplings and the Z' mass in comparison with V_{cb} and the W mass [37]. Consequently, we do not study its effect on ΔM_{B_d} and $\sin 2\beta$ in charmed modes. Nevertheless, it can have significant effects on the $B_s - \bar{B}_s$ system if the couplings $B_{sb}^{L,R}$ are not too small, as we assume in the current analysis.

A. Decay amplitude and branching ratio

In the generalized factorization approach [36], the $\bar{B}_d \rightarrow \phi \bar{K}^0$ decay amplitude is

$$\bar{A}(\bar{B}_d \rightarrow \phi \bar{K}^0) = -\frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* \left[a_3 + a_4 + a_5 - \frac{1}{2} (a_7 + a_9 + a_{10}) \right] X^{(BK, \phi)}, \quad (15)$$

where

$$X^{(BK, \phi)} \equiv \langle \phi | (\bar{s}_\alpha s_\alpha)_{V-A} | 0 \rangle \langle \bar{K} | (\bar{s}_\beta b_\beta)_{V-A} | \bar{B} \rangle = 2f_\phi m_\phi F_1^{BK}(m_\phi^2)(\epsilon^* \cdot p_B) \quad (16)$$

is a factorizable hadronic matrix element. The coefficients a_i are given by

$$a_{2i-1} = c_{2i-1}^{\text{eff}} + \frac{1}{N_C^{\text{eff}}} c_{2i}^{\text{eff}}, \quad a_{2i} = c_{2i}^{\text{eff}} + \frac{1}{N_C^{\text{eff}}} c_{2i-1}^{\text{eff}}, \quad (17)$$

where c_i^{eff} are effective Wilson coefficients that should be used when one replaces the one-loop hadronic matrix elements in the effective Hamiltonian with the corresponding tree-level ones [36]. Nonfactorizable effects are encoded in the effective number of colors N_C^{eff} . Throughout this paper, we take the naive choice $N_C^{\text{eff}} = 3$ for illustration.

For the input parameters $\alpha_s(M_Z) = 0.118$, $\alpha_{\text{EM}} = 1/128$; the Wolfenstein parameters [38] $\lambda = 0.2240$, $A = 0.825$, $\rho = 0.21$ and $\eta = 0.34$ [39]; $\sin^2 \theta_W = 0.23$, $M_W = 80.42$ GeV; and the running quark masses $m_t = 168$ GeV, $m_b = 4.88$ GeV, $m_c = 1.5$ GeV, $m_s = 122$ MeV, $m_u = 4.2$ MeV, and $m_d = 7.6$ MeV [40], the next-to-leading order (NLO) effective Wilson coefficients [35, 36] for the $|\Delta S| = 1$ weak Hamiltonian at the scale $\mu = 2.5$ GeV within the SM are given in the second and third columns of Table I.

The $B_d \rightarrow \phi K^0$ decay width is given by

$$\Gamma(B_d \rightarrow \phi K^0) = \frac{p_c^3}{8\pi m_\phi^2} \left| \frac{A(B_d \rightarrow \phi K^0)}{\epsilon^* \cdot p_B} \right|^2, \quad (18)$$

where

$$p_c = \frac{\sqrt{[m_B^2 - (m_\phi + m_K)^2][m_B^2 - (m_\phi - m_K)^2]}}{2m_B} \quad (19)$$

TABLE II: Experimental results of the CP -averaged branching ratios (quoted in units of 10^{-6}) and CP asymmetries of the $B \rightarrow \phi K$ and $B \rightarrow \eta' K$ decays. References are given in square brackets. The scale factor S (defined in Ref. [13]) is displayed in parentheses when it is larger than 1.

Mode	BaBar	Belle	CLEO	Avg.
$B(B^0 \rightarrow \phi K^0)$	$8.4_{-1.3}^{+1.5} \pm 0.5$ [43]	$9.0_{-1.8}^{+2.2} \pm 0.7$ [44]	$5.4_{-2.7}^{+3.7} \pm 0.7$ [45]	8.3 ± 1.1
$S_{\phi K_S}$	$0.45 \pm 0.43 \pm 0.07$ [4]	$-0.96 \pm 0.50_{-0.11}^{+0.09}$ [4]	-	-0.147 ± 0.697 ($S = 2.11$)
$A_{\phi K_S}$	$0.38 \pm 0.37 \pm 0.12$ [4]	$-0.15 \pm 0.29 \pm 0.07$ [4]	-	0.046 ± 0.256 ($S = 1.08$)
$B(B^0 \rightarrow \eta' K^0)$	$60.6 \pm 5.6 \pm 4.6$ [46]	$68 \pm 10_{-8}^{+9}$ [47]	$89_{-16}^{+18} \pm 9$ [48]	65.18 ± 6.18 ($S = 1.03$)
$S_{\eta' K_S}$	$0.02 \pm 0.34 \pm 0.03$ [4]	$0.43 \pm 0.27 \pm 0.05$ [4]	-	0.269 ± 0.214
$A_{\eta' K_S}$	$-0.10 \pm 0.22 \pm 0.03$ [4]	$-0.01 \pm 0.16 \pm 0.04$ [4]	-	-0.042 ± 0.132

is the momentum of the decay particles in the center-of-mass frame. With $\tau_{B^0} = 1.534$ ps [41], $f_\phi = 237$ MeV, $F_1^{BK}(m_\phi^2) = 0.407$ [42] and meson masses given in Ref. [13], the CP -averaged branching ratio in the SM is

$$\mathcal{B}^{\text{SM}}(B^0 \rightarrow \phi K^0) \simeq 11 \times 10^{-6}. \quad (20)$$

This result is slightly above the 95%CL range of the current world average value $(8.3 \pm 1.1) \times 10^{-6}$ given in Table II, but is close to the previous calculation [16]. (We ignore small theoretical uncertainties in the SM here and in illustrating the consequences of Z' physics in the following sections.)

With FCNC, the Z' boson contributes at tree level, and its contribution will interfere with the standard model contributions. In particular, the flavor-changing couplings of the Z' with the left-handed fermions will contribute to the O_9 and O_7 operators for left (right)-handed couplings at the flavor-conserving vertex, i.e., $c_{9,7}(M_W)$ receive new contributions from Z' . On the other hand, the right-handed flavor changing couplings yield new operators with coefficients that contain another weak phase associated with B_{sb}^R . We will ignore these contributions in this paper.

The effective Hamiltonian of the $b \rightarrow s\bar{s}s$ transition mediated by the Z' is

$$\mathcal{H}_{\text{eff}}^{Z'} = -\frac{4G_F}{\sqrt{2}} \left(\frac{g' M_Z}{g_Y M_{Z'}} \right)^2 B_{sb}^L (B_{ss}^L O_9 + B_{ss}^R O_7) + \text{h.c.}, \quad (21)$$

where $g_Y = e/(\sin\theta_W \cos\theta_W)$, and B_{ij}^L and B_{ij}^R refer to the left- and right-handed effective Z' couplings of the quarks i and j at the weak scale, respectively. The diagonal elements are real due to the hermiticity of the effective Hamiltonian, but the off-diagonal elements may contain weak phases. Only one new weak phase associated with B_{sb}^L can be introduced into the theory under our assumption of neglecting B_{sb}^R . We denote this by ϕ_L and write $B_{sb}^L = |B_{sb}^L| e^{i\phi_L}$. As $\mathcal{H}_{\text{eff}}^{Z'}$ has the same operators O_9 and O_7 as in the SM effective Hamiltonian, the strong phases from long-distance physics should be the same.

Since heavy degrees of freedom in the theory have already been integrated out at the scale of M_W , the RG evolution of the Wilson coefficients after including the new contributions from Z' is exactly the same as in the SM. We obtain the branching ratio

$$\mathcal{B}^{\text{SM}+Z'}(\bar{B} \rightarrow \phi \bar{K}^0) \simeq \mathcal{B}^{\text{SM}}(\bar{B} \rightarrow \phi \bar{K}^0) \left| 1 - [(41.8 - 7.1i)\xi^{LL} + (46.2 - 8.6i)\xi^{LR}] e^{i\phi_L} \right|^2, \quad (22)$$

where

$$\xi^{LL} \equiv \left| \left(\frac{g' M_Z}{g_Y M_{Z'}} \right)^2 \frac{B_{sb}^L B_{ss}^L}{V_{tb} V_{ts}^*} \right|, \quad (23)$$

$$\xi^{LR} \equiv \left| \left(\frac{g' M_Z}{g_Y M_{Z'}} \right)^2 \frac{B_{sb}^L B_{ss}^R}{V_{tb} V_{ts}^*} \right|, \quad (24)$$

and $|V_{tb} V_{ts}^*| \simeq 0.04$. The second and third terms in Eq. (22) represent the Z' contributions from left- and right-handed couplings with the $s\bar{s}$ in the final state, respectively. We have assumed for definiteness that B_{ss}^L and B_{ss}^R have the same sign, so that the ξ^{LL} and ξ^{LR} terms interfere constructively. The branching ratio predicted by our model depends on the absolute ratios ξ^{LL} and ξ^{LR} and the weak phase ϕ_L .

We show the branching ratios as a function of ϕ_L in Fig. 1. Generically, one expects a ratio $g'/g_Y \sim \mathcal{O}(1)$ and $M_{Z'}$ to be a few to around 10 times M_Z . We assume that the product $|B_{sb}^L B_{ss}^L|$ is numerically about the same as $|V_{tb} V_{ts}^*|$,

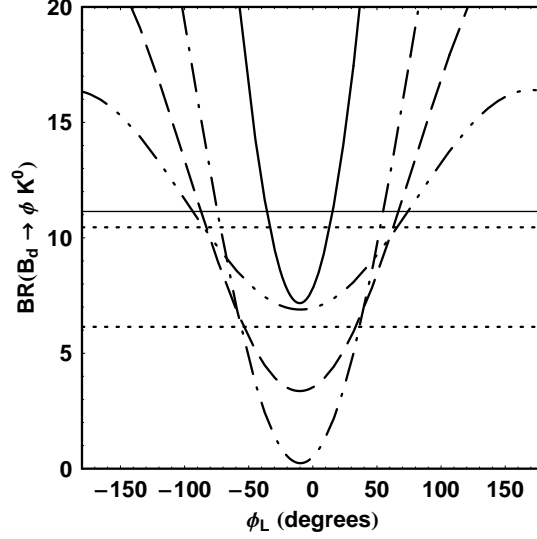


FIG. 1: The branching ratio $\mathcal{B}^{\text{SM}+Z'}(B \rightarrow \phi K^0)$ in units of 10^{-6} versus the weak phase ϕ_L associated with the effective Z' coupling B_{sb}^L . The current experimental range at 95%CL is shown by the two horizontal dotted lines. The SM prediction is the thin horizontal line. The thick solid and dashed curves include both left-handed and right-handed couplings with $\xi^{LL} = \xi^{LR} = 0.02$ and 0.005 , respectively. The single-dot-dashed and double-dot-dashed curves involve only the left-handed couplings with $\xi^{LL} = 0.02$ and 0.005 , respectively.

and take $\xi^{LL} = 0.02$ and 0.005 as representative values for numerical analyses in this and the following sections [60]. It is straightforward to scale the results to other ξ^{LL} values.

To quantify the effects of right-handed couplings, we consider $\xi^{LR} = 0.02$ and 0.005 and show the corresponding curves in Fig. 1. The branching ratio curves are almost symmetric about $\phi_L = 0$, with the slight asymmetry set by the small strong phases in the Wilson coefficients. This echoes the fact that the contributing amplitudes in Eq. (22) have the largest constructive interference when $\phi_L \simeq 0$. To be consistent with the measured branching ratio of $B^0 \rightarrow \phi K^0$, our weak phase ϕ_L in the region $-80^\circ \sim 60^\circ$ is favored, with the exact range depending upon ξ^{LL} and ξ^{LR} in the model. For some parameter choices, it can leave us a two-fold ambiguity, which can be resolved using further information to be discussed in the following subsection.

B. Time-dependent CP asymmetries

The time-dependent CP asymmetry for $B \rightarrow \phi K_S$ is

$$\begin{aligned} a_{\phi K_S}(t) &= \frac{\Gamma(\overline{B}^0(t) \rightarrow \phi K_S) - \Gamma(B^0(t) \rightarrow \phi K_S)}{\Gamma(\overline{B}^0(t) \rightarrow \phi K_S) + \Gamma(B^0(t) \rightarrow \phi K_S)} \\ &= A_{\phi K_S} \cos(\Delta M_{B_d} t) + S_{\phi K_S} \sin(\Delta M_{B_d} t), \end{aligned} \quad (25)$$

where the direct and the indirect CP asymmetry parameters are given respectively by

$$A_{\phi K_S} = \frac{|\lambda_{\phi K_S}|^2 - 1}{|\lambda_{\phi K_S}|^2 + 1}, \quad S_{\phi K_S} = \frac{2\text{Im}[\lambda_{\phi K_S}]}{|\lambda_{\phi K_S}|^2 + 1}. \quad (26)$$

The parameter $\lambda_{\phi K_S}$ is defined by

$$\lambda_{\phi K_S} \equiv \eta_{\phi K_S} \left(\frac{q}{p} \right)_B \left(\frac{p}{q} \right)_K \frac{\overline{A}(\phi \overline{K}^0)}{A(\phi K^0)}, \quad (27)$$

where $\eta_{\phi K_S} = -1$ is the CP eigenvalue of the ϕK_S state, and

$$\left(\frac{q}{p} \right)_B = \frac{V_{tb}^* V_{td}}{V_{tb} V_{td}^*}, \quad \left(\frac{p}{q} \right)_K = \frac{V_{cs} V_{cd}^*}{V_{cs}^* V_{cd}} \quad (28)$$

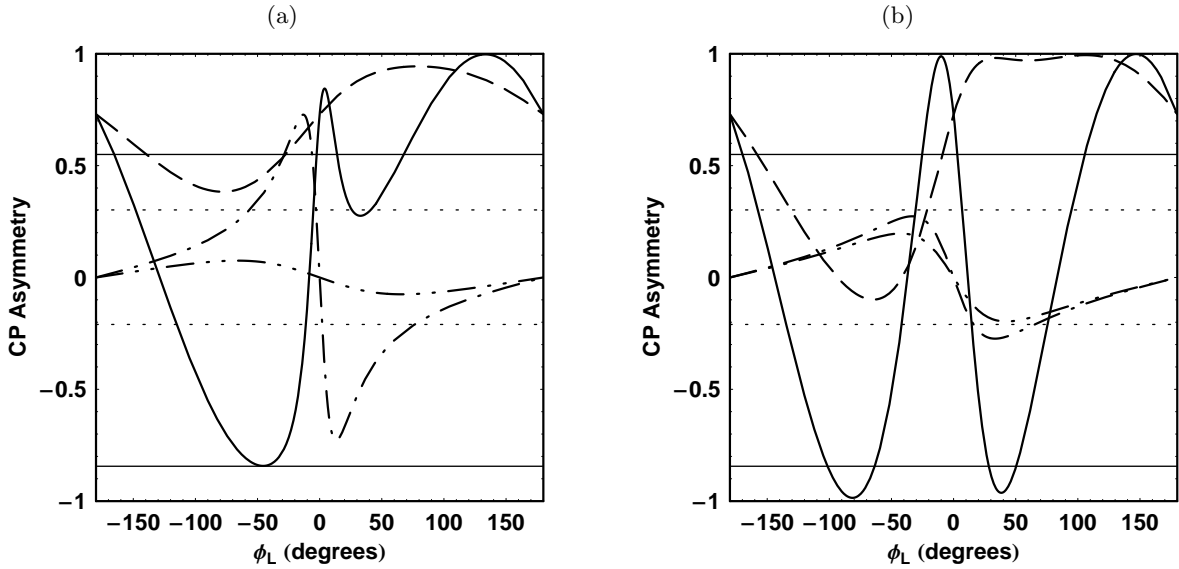


FIG. 2: The time-dependent CP asymmetries, $S_{\phi K_S}$ and $A_{\phi K_S}$, versus ϕ_L . The current experimental ranges at 1σ level are shown by the thin horizontal solid and dotted lines, respectively. The SM predicts $(S_{\phi K_S}, A_{\phi K_S}) \simeq (0.73, 0)$ (not shown). The thick solid and dashed curves are $S_{\phi K_S}$ for $\xi^{LL} = 0.02$ and 0.005 , respectively. The corresponding curves for $A_{\phi K_S}$ are displayed using single-dotted and double-dotted dashes. Plot (a) has $\xi^{LR} = 0$; plot (b) has $\xi^{LR} = \xi^{LL} = 0.02$ and 0.005 .

are factors that account for the mixing effects in neutral B and K meson systems, respectively.

We show our estimates of $S_{\phi K_S}$ and $A_{\phi K_S}$ as a function of the new weak phase ϕ_L in Fig. 2. In Fig. 2(a), we have choices $\xi^{LL} = 0.02$ and 0.005 , but set $\xi^{LR} = 0$. The SM prediction of $S_{\phi K_S}$ and $A_{\phi K_S}$ are 0.73 and 0 , respectively. We see that the measured $A_{\phi K_S}$ does not give much constraint on the weak phase ϕ_L , except for the regions between $-55^\circ \sim 80^\circ$ when $\xi^{LL} = 0.02$. The $S_{\phi K_S}$ data can be readily fitted within 1σ for values of ξ^{LL} chosen here.

In Fig. 2(b), we turn on the right-handed couplings and set $\xi^{LR} = \xi^{LL}$. We notice that the variation of $A_{\phi K_S}$ is within the experimental 1σ limits for the most range of ϕ_L . As illustrated by the thick solid curve in Fig. 2(b), there are four possible ranges of ϕ_L that can fit the averaged $S_{\phi K_S}$ if both ξ^{LL} and ξ^{LR} are large enough. For smaller ξ^{LL} and ξ^{LR} , however, only a region of negative ϕ_L is favored.

In order to satisfy both CP asymmetry constraints, ϕ_L should have negative value in most cases. Only $\xi^{LL} = \xi^{LR} = 0.02$ can have some positive ϕ_L range. Combining the constraints from $\mathcal{B}(B \rightarrow \phi K)$ at 95%CL and both $S_{\phi K_S}$ and $A_{\phi K_S}$ at 1σ level, we find the following allowed regions of ϕ_L . If we take $\xi^{LR} = \xi^{LL} = 0.02$, $5^\circ \lesssim \phi_L \lesssim 15^\circ$ is favored. If we take $\xi^{LR} = \xi^{LL} = 0.005$, $-80^\circ \lesssim \phi_L \lesssim -55^\circ$ is favored. If we take $\xi^{LL} = 0.02$ and ignore ξ^{LR} , then $-70^\circ \lesssim \phi_L \lesssim -55^\circ$ is favored. For $\xi^{LL} = 0.005$ and $\xi^{LR} = 0$, $-80^\circ \lesssim \phi_L \lesssim -30^\circ$ is favored.

IV. $B_d \rightarrow \eta' K_S$

The $B_d \rightarrow \eta' K_S$ is another decay mode whose time-dependent CP asymmetry $S_{\eta' K_S}$ is expected to give us the SM $\sin 2\beta$. Current data reported by BaBar and Belle (see Table II) are both lower than the SM prediction, although consistent within 2σ . Since this process also contains the O_7 and O_9 operators in the amplitude, we discuss the Z' effects on its observables.

The perturbative calculations of the $B \rightarrow \eta' K$ branching ratios are significantly smaller than the observed values. This discrepancy can be explained by adding a singlet-penguin amplitude, where η' is produced through a flavor-singlet neutral current, to interfere constructively with the QCD penguin contributions [49, 50]. Another analysis [51] found that it is hard to obtain a sizeable flavor-singlet amplitude from perturbative calculations, but QCD penguin amplitudes can be enhanced by an asymmetric treatment of the $s\bar{s}$ component of the η' wavefunction. Since this matter is still debatable, we will follow the usual effective Hamiltonian approach [35, 36] and put the emphasis on what kind of effects the Z' boson may provide.

Following the notation in Ref. [16], the decay amplitude of $\bar{B}^0 \rightarrow \eta' \bar{K}^0$ can be written as

$$\bar{A}(\bar{B}^0 \rightarrow \eta' \bar{K}^0)$$

$$\begin{aligned}
&= i \frac{G_F}{\sqrt{2}} \left[V_{ub} V_{us}^* a_2 X_2 - V_{tb} V_{ts}^* \left\{ \left[a_4 - \frac{a_{10}}{2} + \left(a_6 - \frac{a_8}{2} \right) R_1 \right] X_1 + \left[2(a_3 - a_5) - \frac{1}{2}(a_7 - a_9) \right] X_2 \right. \right. \\
&\quad \left. \left. + \left[a_3 + a_4 - a_5 + \frac{1}{2}(a_7 - a_9 - a_{10}) + \left(a_6 - \frac{a_8}{2} \right) R_2 \right] X_3 \right\} \right] , \tag{29}
\end{aligned}$$

where

$$\begin{aligned}
X_1 &= i(m_B^2 - m_{\eta'}^2) \frac{X_{\eta'}}{\sqrt{2}} f_K F_0^{B\pi}(m_{K^0}^2) , \\
X_2 &= i(m_B^2 - m_{K^0}^2) \frac{X_{\eta'}}{\sqrt{2}} f_\pi F_0^{BK}(m_{\eta'}^2) , \\
X_3 &= i(m_B^2 - m_{K^0}^2) Y_{\eta'} \sqrt{2f_K^2 - f_\pi^2} F_0^{BK}(m_{\eta'}^2) , \\
R_1 &= \frac{2m_{K^0}^2}{(m_b - m_d)(m_s + m_d)} , \\
R_2 &= \frac{2(2m_{K^0}^2 - m_\pi^2)}{2m_s(m_b - m_s)} , \tag{30}
\end{aligned}$$

and $X_{\eta'} = 0.57$ and $Y_{\eta'} = 0.82$ are mixing parameters for the choice of the η' meson wavefunction to be $(2s\bar{s} + u\bar{u} + d\bar{d})/\sqrt{6}$.

Since the $B \rightarrow \eta' K$ has two pseudoscalar mesons in the final state, the decay width is

$$\Gamma(B_d \rightarrow \eta' K^0) = \frac{p_c}{8\pi m_B^2} |A(B_d \rightarrow \eta' K^0)|^2 , \tag{31}$$

where p_c is defined in a similar way to Eq. (19). With $f_\pi = 131$ MeV and $f_K = 159.8$ MeV [13], $F_0^{B\pi}(m_{K^0}^2) = 0.335$ and $F_0^{BK}(m_{\eta'}^2) = 0.391$ [42], we have $\mathcal{B}^{\text{SM}}(B \rightarrow \eta' K^0) \simeq 38 \times 10^{-6}$, which is much lower than the experimental average of $(65.18 \pm 6.18) \times 10^{-6}$ (see Table II).

As in the case of $B \rightarrow \phi K$ decays, our model makes extra contributions to O_9 and O_7 at the weak scale. The branching ratio is

$$\mathcal{B}^{\text{SM}+Z'}(\bar{B} \rightarrow \eta' \bar{K}^0) \simeq \mathcal{B}^{\text{SM}}(\bar{B} \rightarrow \eta' \bar{K}^0) \left| 1 - [(7.0 - 0.5i)\xi^{LL} + (2.9 - 0.4i)\xi^{LR}] e^{i\phi_L} \right|^2 . \tag{32}$$

We notice that the coefficient of ξ^{LL} and that of ξ^{LR} also tend to have constructive interference between themselves according to our assumption that B_{ss}^L and B_{ss}^R have the same sign. The magnitudes of these coefficients, however, are much smaller than those in Eq. (22). This is simply because the terms that receive contributions from the Z' boson (mostly a_9) have some cancellation between the X_2 and X_3 terms in Eq. (29). These observations qualitatively tell us why the $\eta' K$ decays are not affected quite as much by the Z' effects.

We see in Fig. 3 that the Z' boson can explain the gap between the observed branching ratio and the SM prediction only around $\phi_L = \pm 180^\circ$ even with large couplings in both ξ^{LL} and ξ^{LR} . As we will see, however, this region is not favored by the CP asymmetry constraints. Therefore, we must attribute this anomaly to some other unknown source.

The asymmetry curves for $B_d \rightarrow \eta' K_S$ are shown in Fig. 4. We do not get useful constraints from current data on $A_{\eta' K_S}$. The value of $A_{\eta' K_S}$ does not vary much from its SM prediction throughout the whole range of ϕ_L . The averaged value of $S_{\eta' K_S}$ can be explained at 1σ level by simultaneously taking large values of both left- and right-handed couplings (the solid curve in Fig. 4(b)). In this case, however, only negative ϕ_L around $-120^\circ \sim -40^\circ$ is favored from the $S_{\eta' K_S}$ constraint. Other cases do not explain the $S_{\eta' K_S}$ anomaly though all of them favor negative value of ϕ_L to approach the 1σ limit.

Leaving 95%CL of branching ratio constraints, we have only $\xi^{LL} = \xi^{LR} = 0.02$ case that can satisfy both $B \rightarrow \eta' K$ and $B \rightarrow \eta' K_S$ CP asymmetry with a two-fold range of ϕ_L , $-120^\circ \sim -100^\circ$ and $-60^\circ \sim -40^\circ$. Attributing the branching ratio of $B \rightarrow \eta' K$ to some unknown effects, the latter is favored by the $B \rightarrow \phi K$ branching ratio.

V. $B_s \rightarrow \mu^+ \mu^-$

$B \rightarrow \ell^+ \ell^-$ decays are good candidates to observe FCNC interactions beyond the SM. In the SM, the branching ratio of such processes are proportional to m_ℓ^2 and the corresponding CKM factors. Therefore, B_s decays in general have branching ratios larger than the corresponding B_d decays by a factor of $\sim \lambda^2$. Although $B_s \rightarrow \tau^+ \tau^-$ should

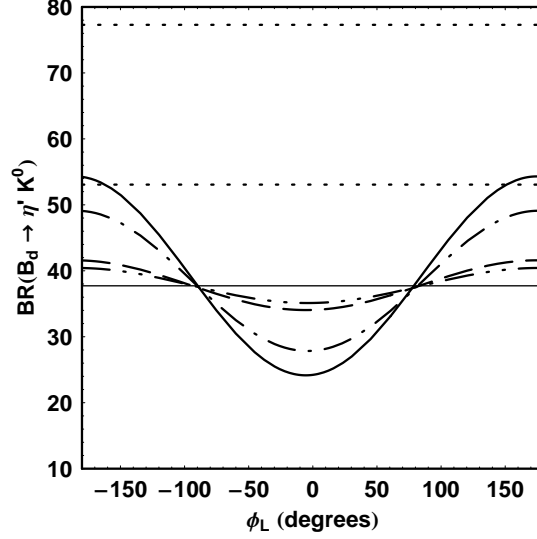


FIG. 3: The branching ratio $\mathcal{B}^{\text{SM}+Z'}(B \rightarrow \eta' K^0)$ in units of 10^{-6} versus ϕ_L . The current experimental range at 95%CL is shown by the two horizontal dotted lines. The SM prediction is the thin horizontal line. The thick solid and dashed curves include both left-handed and right-handed couplings with $\xi^{LL} = \xi^{LR} = 0.02$ and 0.005 , respectively. The single-dot-dashed and double-dot-dashed curves involve only the left-handed couplings with $\xi^{LL} = 0.02$ and 0.005 , respectively.

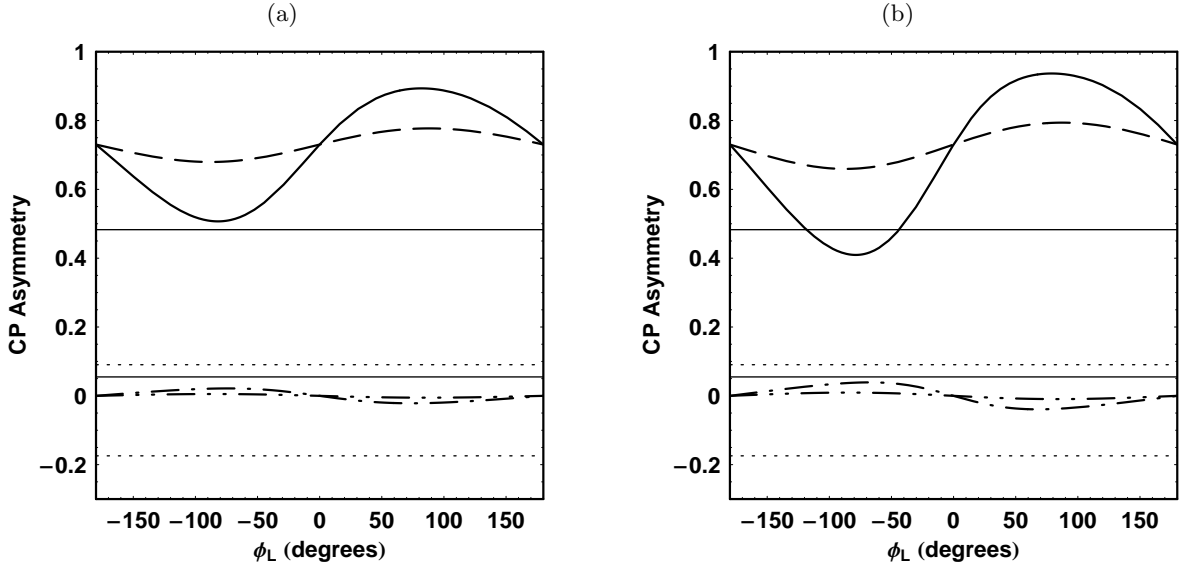


FIG. 4: The time-dependent CP asymmetries, $S_{\eta'K_S}$ and $A_{\eta'K_S}$, versus ϕ_L . The current experimental ranges at 1σ level are shown by the thin horizontal solid and dotted lines, respectively. The SM predicts $(S_{\eta'K_S}, A_{\eta'K_S}) \simeq (0.73, 0)$ (not shown). The thick solid and dash-dotted curves are $S_{\eta'K_S}$ for $\xi^{LL} = 0.02$ and 0.005 , respectively. The corresponding curves for $A_{\eta'K_S}$ are displayed using single-dotted and double-dotted dashes. Plot (a) has $\xi^{LR} = 0$; plot (b) has $\xi^{LR} = \xi^{LL} = 0.02$ and 0.005 .

have a still larger branching ratio, this mode is hard to detect at hadron colliders. The best mode for experimental study is thus $B_s \rightarrow \mu^+ \mu^-$.

Recently, studies in some minimal supersymmetric standard model (MSSM) and minimal supergravity (mSUGRA) models predict that the branching ratio of the rare decay $B_s \rightarrow \mu^+ \mu^-$ can be large enough to be observable by Tevatron Run II [52]. Using 113 pb^{-1} of data from Run II, the CDF collaboration placed an upper bound $\mathcal{B}(B_s \rightarrow \mu^+ \mu^-) < 9.5 \times 10^{-7}$, while the D0 collaboration placed a bound $< 16 \times 10^{-7}$ based on 100 pb^{-1} of data [53]. The SM prediction of $\mathcal{B}(B_s \rightarrow \mu^+ \mu^-)$ is more than two orders of magnitude smaller than these bounds.

The SM contribution to $B_s \rightarrow \mu^+ \mu^-$ is loop-suppressed. It is therefore possible for the decay to be dominated by Z' physics. For an order of magnitude estimate, one can temporarily ignore the RG running effect at the b - s - Z'

vertex. The decay width of $B_s \rightarrow \mu^+ \mu^-$ is given by

$$\mathcal{B}(B_s \rightarrow \mu^+ \mu^-) = \tau(B_s) \frac{G_F^2}{4\pi} f_{B_s}^2 m_\mu^2 m_{B_s} \sqrt{1 - \frac{4m_\mu^2}{m_{B_s}^2}} |V_{tb}^* V_{ts}|^2 \times \left\{ \left| \frac{\alpha}{2\pi \sin^2 \theta_W} Y \left(\frac{m_t^2}{M_W^2} \right) + 2 \left(\frac{g' M_Z}{g_Y M_{Z'}} \right)^2 \frac{B_{bs}^L B_{\mu\mu}^L}{V_{tb}^* V_{ts}} \right|^2 + \left| 2 \left(\frac{g' M_Z}{g_Y M_{Z'}} \right)^2 \frac{B_{bs}^L B_{\mu\mu}^R}{V_{tb}^* V_{ts}} \right|^2 \right\}, \quad (33)$$

where $B_{\mu\mu}^{L,R}$ is the effective μ - μ - Z' coupling at the weak scale. One can find the definition of $Y(m_t^2/M_W^2)$ in the SM part in Ref. [54]; its value is about 1.05 here. No mixing between Z and Z' is assumed here. Using the central value of the averaged B_s lifetime $\tau_{B_s} = 1.439$ ps [41] and $f_{B_s} = 232$ MeV, we obtain a SM branching ratio of $\simeq 4.1 \times 10^{-9}$. The current CDF Run II upper bound thus gives a constraint that

$$\left\{ \left| 0.0028 + \left(\frac{g' M_Z}{g_Y M_{Z'}} \right)^2 \frac{B_{bs}^L B_{\mu\mu}^L}{V_{tb}^* V_{ts}} \right|^2 + \left| \left(\frac{g' M_Z}{g_Y M_{Z'}} \right)^2 \frac{B_{bs}^L B_{\mu\mu}^R}{V_{tb}^* V_{ts}} \right|^2 \right\}^{1/2} \lesssim 0.043. \quad (34)$$

The weak phase ϕ_L associated with B_{sb}^L can in principle be extracted from the time-dependent CP asymmetry measurement in a fashion similar to the cases studied in the previous sections, but the relevant experimental information is not presently available.

We note that the above bound has no direct relation with the couplings relevant to $b \rightarrow s\bar{s}s$ transitions. The lepton couplings to Z' can be much smaller than the quark couplings, as is true in some (quasi)leptophobic models. We are currently investigating the Z' effects on $b \rightarrow s\ell^+\ell^-$ decays, which have been measured to good precision recently and should provide a tighter bound [55, 56].

VI. CONCLUSIONS

In this paper we have considered models with an extra Z' in the mass range of a few hundred GeV to around 1 TeV. With a family nonuniversal structure in the Z' couplings, flavor changing neutral currents are induced via the fermion mixing, therefore producing interesting effects. Currently, constraints on the Z' coupling between the second and third generations are not restrictive. With non-diagonal left-handed and diagonal right-handed Z' couplings in the down-type quarks, we studied the impact of such Z' models on rare B meson decay processes that are sensitive to new physics.

In the present analysis, we have assumed that the left- and right-chiral couplings B_{ss}^L and B_{ss}^R have the same sign, rendering constructive interference in the Z' contributions. We do not include the right-handed flavor changing couplings, which will give rise to new operators not existent in the SM. Involving these or choosing different values for the effective number of colors N_C^{eff} , for which the branching ratios change sensitively, would change the results.

We have found that with the inclusion of the Z' contributions, $S_{\phi K_S}$ can be appreciably different from the SM prediction, while the branching ratio of $B^0 \rightarrow \phi K^0$ and $A_{\phi K_S}$ are still within the experimental ranges. We find that a sizeable weak phase associated with the B_{sb}^L coupling is favored in the ranges of $-80^\circ \sim -30^\circ$, depending upon the ξ^{LL} and ξ^{LR} parameter choices.

We have also studied the influence of the new Z' on the $B^0 \rightarrow \eta' K^0$ decay. The $A_{\eta' K_S}$ data do not restrict the choice of ϕ_L . The $S_{\eta' K_S}$ constraint from the data can be satisfied if large couplings are taken. Though the discrepancy between the observed branching ratio and the SM prediction can be explained with this Z' effect, we cannot explain both branching ratio and CP asymmetries constraints with a common weak angle. Combining with the constraints from the $B^0 \rightarrow \phi K^0$ decays, $S_{\eta' K_S}$ and $A_{\eta' K_S}$, we find that a value of ϕ_L around $-60^\circ \sim -40^\circ$ is favored.

We have observed that the CP asymmetries of the ϕK_S mode are more sensitive to the Z' effects than the $\eta' K_S$ decay. This is because of a cancellation between different parts ($s\bar{s}$ versus $u\bar{u}$ and $d\bar{d}$) in the η' wavefunction.

Finally, we have investigated $B_s \rightarrow \mu^+ \mu^-$ decay, now being searched for at Fermilab, in the same Z' model. This process can be dominated by the Z' contribution and the branching ratio can reach the expected sensitivity in Run II.

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- [57] see Ref. [13] for the definition of the scale factor.
- [58] Such mixings would modify the expressions for the phenomenological ξ^{LL} and ξ^{LR} parameters defined in Section III A, but would not alter the discussion of the implications.
- [59] One could instead always work in the $V_{uL} = \mathbf{1}$ basis, in which case $\epsilon_{dL,R}$ would in general be off-diagonal and complex.
- [60] One may want to compare our choices of ξ^{LL} here with the bounds on the product of the effective s - b - Z' coupling and the ratio $M_Z^2/M_{Z'}^2$ obtained from certain semileptonic B decays in Ref. [28], although a specific E_6 model with leptophobic features is assumed in their study.